

Natural Higgs-flavor-democracy solution of the μ problem of supersymmetry and the QCD axion

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We show that the hierarchically small μ term in supersymmetric theories is a consequence of two identical pairs of Higgs doublets taking a democratic form for their mass matrix. We briefly discuss the discrete symmetry $S_2 \times S_2$ toward the democratic mass matrix. Then, we show that there results an approximate Peccei-Quinn symmetry and hence the value μ is related to the axion decay constant.

PACS numbers: 12.60.Jv, 11.30.Er, 14.80.Va, 14.80.Da

Keywords: μ problem, Approximate PQ symmetry, QCD axion

I. INTRODUCTION

The large hadron collider (LHC) at CERN in Geneva, Switzerland is designed to study the mass scale problem in physics. The Standard Model (SM) of particle physics describes most of the observed phenomena in the universe with the W and Z boson mass scale of order 100 GeV. For the W and Z boson masses and also for the quark and lepton masses [1], the Higgs field seems responsible. The two humongous multi-purpose detectors, CMS and ATLAS, at CERN seem to support the existence of the SM Higgs boson through its Yukawa couplings to the quarks [2]. For the W and Z boson masses and the quark and lepton masses, spontaneous breaking of the SM gauge symmetry is the accepted possibility by the vacuum expectation value (VEV) of the so-called Higgs doublet $H = \{H^0, H^-\}$.

Supersymmetry is a mild extension of the SM. The simplest extension the Minimal Supersymmetric Standard Model (MSSM) introduces a superpartner for every SM particle except for the Higgs doublet. The Higgs doublet is extended to one pair $H_u = \{H_u^+, H_u^0\}$ and $H_d = \{H_d^0, H_d^-\}$ and then supersymmetrized. In the MSSM, however, there are two serious problems: the μ problem [3] and the strong CP problem [4]. The μ term, $\mu H_u H_d$, gives order μ Higgs boson mass. The problem is that this μ term breaks no low energy symmetry and is naturally expected to be of order the Planck scale $M_P \simeq 2.44 \times 10^{18}$ GeV or the grand unification (GUT) scale $M_G \simeq 2.5 \times 10^{16}$ GeV.¹ These hierarchically large mass scales are shown in the first column of Fig. 1. The strong CP problem is a problem on the QCD vacuum angle θ which is bounded to the extremely small region, $|\theta| < 10^{-10}$ [8]. The invisible axion solution [9] of the strong CP problem is based on the Peccei-Quinn (PQ) global symmetry [4]. However, the PQ global symmetry

is not warranted in gravitational interactions [10]. Even though the recent LHC data [2] seems to support the existence of the SM Higgs boson but no superpartner of the SM has been found so far below the TeV scale. Thus, the 100 GeV scale SUSY scenario in the family universal MSSM is in doubt but it is too early to discard a TeV scale MSSM altogether. For example, the family dependent U(1) quantum numbers can save the MSSM [11], so we continue to take the supersymmetry solution of the gauge hierarchy problem here.

In order to address these problems with supersymmetry, we will work only with discrete symmetries in this Letter so that the gravity argument against axion [10] does not apply here. Nevertheless, an approximate PQ symmetry will be derived from the discrete symmetry and solves the strong CP problem. In this way, the axion decay constant is related to the supersymmetry parameter μ .

From the very high energy point of view, the *fundamental problem* of μ is “Why is the parameter μ in the MSSM so small compared to the GUT scale?” If this problem is going to be understood from a symmetry principle, the first step toward the solution is obtaining the zero value of μ at the lowest order as depicted in the second column of Fig. 1. The next important question on μ is: “Why does the parameter μ fall in the electroweak scale range?”

Obtaining massless particles naturally was considered long time ago for three families of fermions right after the discovery of τ and b under the name of ‘flavor democracy’ [14]. The chief motivation of Ref. [14] was in order to obtain the heavy third family fermions, rather than to obtain two massless fermions of the first two families. For the flavor democracy of three up-type quarks, one introduces the permutation symmetries on three flavors, independently for the left-handed(L) fields and for the right-handed(R) fields: $S_3(L) \times S_3(R)$. Then, the mass matrix for the up-type quarks takes the form,

$$\begin{pmatrix} m_t/3 & m_t/3 & m_t/3 \\ m_t/3 & m_t/3 & m_t/3 \\ m_t/3 & m_t/3 & m_t/3 \end{pmatrix},$$

¹ We note that the μ problem has been discussed long time ago with the PQ symmetry [3], with supergravity effects [5] and in string models [6]. More recently, the discrete Z_4^R symmetry has been discussed also [7].

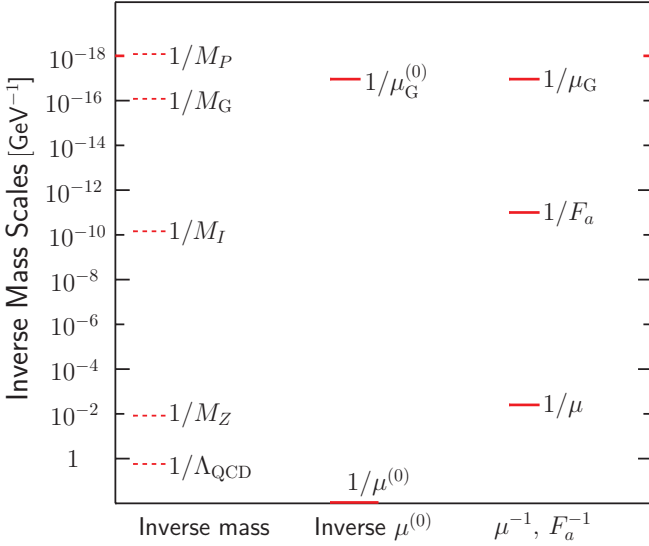


FIG. 1: Interaction scales. The scale $\mu^{(0)}$ is for the lowest order toward a natural solution, and the scale μ is generated at higher orders. Here, M_I is the intermediate scale $M_I = \sqrt{M_P M_Z}$, and the PQ symmetry breaking scale F_a is near M_I .

leading to the mass eigenvalues of m_t , 0, and 0. This kind of $S_3 \times S_3$ is not helpful for obtaining just one pair of massless Higgs doublets, or by supersymmetry one pair of massless Higgsino doublets.

Therefore, to obtain just one massless pair of Higgsino doublets in supersymmetric models, we need $S_2(H_u) \times S_2(\tilde{H}_d)$ for two Higgsino doublet pairs, $\{\tilde{H}_u^{(i)}, \tilde{H}_d^{(i)}\}$ with $i = 1, 2$. Here, $S_2(\tilde{H}_u)$ is the interchange symmetry ($1 \leftrightarrow 2$) for two Higgsino doublets $\tilde{H}_u^{(1)}$ and $\tilde{H}_u^{(2)}$. So is $S_2(\tilde{H}_d)$ for two Higgsino doublets $\tilde{H}_d^{(1)}$ and $\tilde{H}_d^{(2)}$. Due to supersymmetry, we identify the Higgsino discrete symmetry $S_2(\tilde{H}_u) \times S_2(\tilde{H}_d)$ with the Higgs boson discrete symmetry $S_2(H_u) \times S_2(H_d)$, shortened as *Higgs-flavor-democracy*. If we start with just one pair of Higgs doublets, we cannot understand the *fundamental problem* of μ naturally.

II. HIGGS-FLAVOR-DEMOCRACY

Our solution of the μ problem is by introducing just two pairs of Higgsino doublets at the GUT scale M_G . But, the unification of gauge coupling constants at a GUT scale ($M_G \sim 2 \times 10^{16}$ GeV) requires only one pair of Higgs doublets $\{H_u, H_d\}$ at low energy [12], which is understood as in Fig. 1 with just one pair surviving below the GUT scale. If the two Higgsino pairs are not distinguished by any quantum number and geometry of the internal space, there must be the permutation symmetries, $S_2(H_u)$ and $S_2(H_d)$. Then, the democratic form

for the Higgsino mass matrix can be obtained,

$$\begin{pmatrix} M_G/2, & M_G/2 \\ M_G/2, & M_G/2 \end{pmatrix} \quad (1)$$

which gives Higgsino mass eigenvalues M_G and 0. The *Higgs-flavor-democracy* based on $S_2(H_u) \times S_2(H_d)$ works as follows. The representation content of S_2 is only a singlet **1**. The tensor product (for the mass matrix) of two S_2 singlets, *i.e.* $H_u = (H_u^{(1)}, H_u^{(2)})^T$ and $H_d = (H_d^{(1)}, H_d^{(2)})^T$, contains two independent real parameters m_E and m_O for the diagonal and off-diagonal combinations, respectively,

$$W = \left[m_E (H_u^{(1)\alpha} H_d^{(1)\beta} + H_u^{(2)\alpha} H_d^{(2)\beta}) + m_O (H_u^{(1)\alpha} H_d^{(2)\beta} + H_u^{(2)\alpha} H_d^{(1)\beta}) \right] \epsilon_{\alpha\beta} \quad (2)$$

where α and β are the $SU(2)_W$ gauge group indices of the SM. Now, let us apply $S_2(H_u)$ symmetry to W : $1 \rightarrow 2, 2 \rightarrow 1$ for $H_u^{(i)}$, and $1 \rightarrow 1, 2 \rightarrow 2$ for $H_d^{(i)}$. Then, we obtain

$$W \rightarrow \left[m_E (H_u^{(2)\alpha} H_d^{(1)\beta} + H_u^{(1)\alpha} H_d^{(2)\beta}) + m_O (H_u^{(2)\alpha} H_d^{(2)\beta} + H_u^{(1)\alpha} H_d^{(1)\beta}) \right] \epsilon_{\alpha\beta}. \quad (3)$$

Comparing Eqs. (2) and (3), we obtain $m_E = m_O$ and obtain the *Higgs-flavor-democracy*, Eq. (1). Applying $S_2(H_d)$ gives the same result. But if we apply $S_2(H_u)$ and $S_2(H_d)$ simultaneously, m_E and m_O are not related.

Thus, the $S_2(H_u) \times S_2(H_d)$ invariant superpotential is

$$W_{S_2 \times S_2} = \frac{M_G}{2} \left(H_u^{(1)\alpha} H_d^{(1)\beta} + H_u^{(2)\alpha} H_d^{(2)\beta} + H_u^{(1)\alpha} H_d^{(2)\beta} + H_u^{(2)\alpha} H_d^{(1)\beta} \right) \epsilon_{\alpha\beta}. \quad (4)$$

The mass matrix is diagonalized to the new $(H^{(0)}, H^{(G)})^T$ basis,

$$M_0^{\text{Higgsino}} = \begin{pmatrix} 0 & 0 \\ 0 & M_G \end{pmatrix}, \quad H_{u,d}^{(0),(M_G)} = \frac{1}{\sqrt{2}} \left(H_{u,d}^{(1)} \mp H_{u,d}^{(2)} \right). \quad (5)$$

Indeed, one can find a few string models allowing two identical pairs of Higgs doublets in the MSSM [15, 16]. We note that Ref. [16] contains two pairs of Higgs doublets in the twisted sector T_6 , has the *Higgs-flavor-democracy*, and so naturally contains a light pair of Higgs doublets.

III. GENERATION OF TEV SCALE μ

Since the *Higgs-flavor-democracy* gives one pair of the Higgsino doublets zero mass as in the second column of

Fig. 1, one has to break the *Higgs-flavor-democracy* to obtain a TeV scale μ , or the massless Higgsino can never obtain mass. In the supersymmetric field theory framework, we show a possibility that the *Higgs-flavor-democracy* is broken. Let us take the minimal Kähler potential $K = \Phi_i \Phi_i^\dagger$ where Φ_i ($i = 1, 2$) is the gauge group non-singlet field such as the Higgs superfield and X_i ($i = 1, 2$) and \bar{X}_i ($i = 1, 2$) are gauge group singlet superfields, obeying the common $S_2 \times S_2$ symmetry of Φ_i ($i = 1, 2$) and $\bar{\Phi}_i$ ($i = 1, 2$),

$$S_2 : \Phi_1 \leftrightarrow \Phi_2, X_1 \leftrightarrow X_2, \quad (6)$$

and similarly for the barred fields. Let us consider the following $S_2 \times S_2$ symmetric non-renormalizable term [3],

$$W^{(\text{nonren.})} = \sum_{i,j=1,2} \left(\frac{X^{(i)} \bar{X}^{(j)}}{M_P} \right) H_u^{(i)} H_d^{(j)}. \quad (7)$$

To have VEVs of X and \bar{X} fields, let us consider an $S_2 \times S_2$ symmetric superpotential with a singlet Z ($Z \rightarrow Z$ under S_2) [17],

$$W = \lambda Z (X_1 \bar{X}_1 + X_1 \bar{X}_2 + X_2 \bar{X}_1 + X_2 \bar{X}_2 - F_a^2). \quad (8)$$

Here, we removed the tadpole term of the X_i and \bar{X}_i fields by an appropriate matter parity such as $P(X_i) = P(\bar{X}_i) = -1$, and others with the even parity. There exists a flavor-democracy breaking minimum, $\langle Z \rangle = 0$, $\langle X_1 \rangle = \langle \bar{X}_1 \rangle = F_a$, $\langle X_2 \rangle = \langle \bar{X}_2 \rangle = 0$. Since there also exists the $S_2 \times S_2$ symmetric vacuum $\langle Z \rangle = 0$, $X_1 = \bar{X}_1 = X_2 = \bar{X}_2 \neq 0$, our choice of democracy breaking minimum is spontaneous. At the democracy breaking vacuum, $\langle X_1 \rangle = \langle \bar{X}_1 \rangle = F_a$ and $\langle X_2 \rangle = \langle \bar{X}_2 \rangle = 0$, we generate the following term

$$W^{(\text{nonren.})} = \frac{X_1 \bar{X}_1}{2M_P} (H_u^{(0)} + H_u^{(M_G)})(H_d^{(0)} + H_d^{(M_G)}). \quad (9)$$

From Eqs. (5) and (9), we obtain the following Higgsino mass matrix

$$M^{(\text{Higgsino})} = \begin{pmatrix} \mu & \mu \\ \mu & M_G + \mu \end{pmatrix} \quad (10)$$

where $\mu = F_a^2/2M_P$. The eigenvalues of $M^{(\text{Higgsino})}$ are $\mu - \frac{\mu^2}{M_G}$ and $M_G + \mu(1 + \frac{\mu}{M_G})$. Choosing F_a at the intermediate scale $\sim 10^{10-12}$ GeV, we obtain the TeV scale μ .

The common discrete symmetry (6) can be realized in a unification model of Φ_i and X_i . For example, in the SU(6) GUT model [18] H_d is assigned in a $\bar{\mathbf{6}}$: $(T_{r1}^*, T_{g1}^*, T_{b1}^*, H_{d1}^0, -H_{d1}^-, X_1)^T$ and $(T_{r2}^*, T_{g2}^*, T_{b2}^*, H_{d2}^0, -H_{d2}^-, X_2)^T$ where T 's are the color triplet states. In string models, this argument may apply if Φ_i and X_i appear exactly in the same way at a fixed point except for their gauge charges.

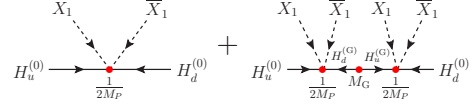


FIG. 2: Diagrams generating the μ term.

IV. INVISIBLE AXION

In the $S_2 \times S_2$ breaking vacuum, Eq. (9), integrating out the heavy fields $H_{u,d}^{(M_G)}$, we have the light field interaction

$$W = \frac{X_1 \bar{X}_1}{2M_P} H_u^{(0)} H_d^{(0)}. \quad (11)$$

The Higgs multiplets $H_u^{(0)}$ and $H_d^{(0)}$ couple to quarks, $W = -qu^c H_u^{(0)} - qd^c H_d^{(0)}$, and define their PQ charges. Then, the PQ charges of X_1 and \bar{X}_1 are given through Eq. (11). This is the appearance of an effective PQ symmetry at low energy from the original discrete symmetry $S_2(H_u) \times S_2(H_d)$. But, this PQ symmetry is not exact if one considers the full Lagrangian including the $H_{u,d}^{(M_G)}$ fields. The related diagrams are shown in Fig. 2. The left diagram of Fig. 2 is the μ term definition consistent with the PQ symmetry. The right diagram breaks the PQ symmetry, gives a correction to μ , which is smaller than the left one by a factor μ/M_G . Thus, the PQ symmetry is approximate, and the explicit PQ symmetry breaking term considering $H_{u,d}^{(M_G)}$ will lead to a very small θ_{QCD} term at the order $\mu/M_G \sim 10^{-14}$, well below the current bound of 10^{-10} [8]. Note that the VEV of X_1 (and \bar{X}_1) is the common scale for breaking the S_2 symmetry toward a nonzero μ and the axion decay constant F_a , anticipated long time ago in [17], to realize the invisible axion [9]. Probably, it is a good rationale that the invisible axion scale is the intermediate scale M_I . So, the axion mass is in the range $10 \mu\text{eV} - 1 \text{meV}$ [8]. It has been observed that gravity is not respecting the PQ global symmetry [10], and hence the components of the antisymmetric tensor field B_{MN} in string models [19] and approximate global symmetries from string [20] were tried for the QCD axion. So, the QCD axion based on the discrete symmetry even at the field theory level is circumventing all these worries. The effective PQ symmetry we obtain here from the matter field X_1 and \bar{X}_1 leads to a reasonable QCD axion from string.

V. ON THE GUTS

Finally, we comment on how the solution of the μ problem is resolved in GUTs. For a complete GUT example, we refer to the flipped SU(5) GUTs [21] from the \mathbb{Z}_{12-I} orbifold compactification [22, 23] and from the fermionic string [24]. In particular, we find two pairs of Higgs doublets in the twisted sector T_4^0 of Ref. [23], as shown in

Sector	Weight	Mult.	$SU(5)_X$
T_4^0	$(\underline{10000}; \frac{1}{3} \frac{1}{3} \frac{1}{3}) (0^8)'$	2	$\mathbf{5}_{-2}(h_u^{(i)})$
T_4^0	$(\underline{-10000}; \frac{1}{3} \frac{1}{3} \frac{1}{3}) (0^8)'$	2	$\mathbf{\bar{5}}_2(h_d^{(i)})$
T_4^0	$(0^5; \frac{-2}{3} \frac{-2}{3} \frac{-2}{3}) (0^8)'$	3	$\mathbf{1}_0(X_i)$
T_3	$(\frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{-1}{2} \frac{-1}{2}; 0^3) (0^5; \frac{-1}{4} \frac{-1}{4} \frac{2}{4})'$	1	$\mathbf{\bar{10}}_{-1}(\bar{H})$
T_9	$(\frac{1}{2} \frac{1}{2} \frac{-1}{2} \frac{-1}{2} \frac{-1}{2}; 0^3) (0^5; \frac{1}{4} \frac{1}{4} \frac{-2}{4})'$	1	$\mathbf{10}_1(H)$
T_2^0	$(0^5; \frac{-1}{3} \frac{-1}{3} \frac{-1}{3}) (0^5; \frac{-1}{2} \frac{1}{2} 0)'$	1 + 1	$\mathbf{1}_0$
T_2^0	$(0^5; \frac{-1}{3} \frac{-1}{3} \frac{-1}{3}) (0^5; \frac{1}{2} \frac{-1}{2} 0)'$	1 + 1	$\mathbf{1}_0$

TABLE I: The T_4^0 , T_3 , T_9 , and T_2^0 left-handed states discussed in the text. The flipped $SU(5)$ quantum numbers are those of $SU(5) \times U(1)_X$.

Table I. But we present the discussion at the supersymmetric field theory level below. The GUT Higgs multiplets $h_u^{(i)} \equiv \mathbf{5}_{-2}^{(i)}$ and $h_d^{(i)} \equiv \mathbf{\bar{5}}_2^{(i)}$ have the infamous doublet-triplet splitting problem. In the flipped $SU(5)$, it is resolved by coupling the color (anti-)triplets in h_u and h_d to the color (anti-)triplets in $\bar{H} \equiv \mathbf{\bar{10}}_{-1}$ from T_3 and $H \equiv \mathbf{10}_1$ from T_9 by $W \sim \sum_i H H h_u^{(i)} + \sum_i \bar{H} \bar{H} h_d^{(i)}$ [24].² But one combination out of two color triplet pairs of h_u and h_d remains massless if $S_2(h_u) \times S_2(h_d)$ is unbroken with the same reason for the case of *Higgs-flavor-democracy*. Including the color (anti-)triplets in \bar{H} and H and color (anti-)triplets in two sets of $h_u^{(i)}$ and $h_d^{(i)}$, we expect the following mass matrix for the color triplets,

$$M_{\text{color}} \propto \begin{pmatrix} \delta & \xi_1 & \xi_2 \\ \eta_1 & M_G & M_G \\ \eta_2 & M_G & M_G \end{pmatrix}$$

where δ, ξ_i and η_i are in general of order M_G . If $S_2(h_u) \times S_2(h_d)$ remains unbroken, we have $\xi_1 = \xi_2$ and $\eta_1 =$

η_2 and $\text{Det. } M_{\text{color}} = 0$, implying one massless pair of color triplet and anti-triplet. For the $S_2(h_u)$ symmetry, we can break it by $W_{h_u} = \lambda_1 H H h_u^{(1)} + \lambda_2 H H h_u^{(2)}$ with $\lambda_1 \neq \lambda_2$, and similarly for the $S_2(h_d)$ symmetry by W_{h_d} . With λ_1 and λ_2 of order 1, all color triplets and anti-triplets become superheavy. In the flipped $SU(5)$, one cannot break the *Higgs-flavor-democracy* for the Higgs doublets via W_{h_u} and W_{h_d} because $\langle H \rangle$ and $\langle \bar{H} \rangle$ do not give mass to the Higgs doublets [24]. Thus, introducing W_{h_u} and W_{h_d} , we achieve the doublet-triplet splitting in the flipped $SU(5)$ GUT.

VI. CONCLUSION

Introducing a global symmetry in string models toward the strong CP solution by spontaneous breaking of the PQ symmetry has been a dilemma for a long time. In this paper, we have found a mechanism to introduce it approximately on the way to solve the μ problem with *Higgs-flavor-democracy*. The underlying symmetry is the discrete $S_2(H_u) \times S_2(H_d)$ symmetry for two identical pairs of Higgs doublets, $\{H_u^{(i)}, H_d^{(i)}\}$ ($i = 1, 2$), at the high energy scale. Being discrete, the $S_2(H_u) \times S_2(H_d)$ symmetry can be realized in string models. In sum, it has not escaped our attention that two identical pairs of Higgs doublets with supersymmetry introduce *Higgs-flavor-democracy*, bring one pair down to the TeV scale solving the μ problem naturally, and predicts a very light axion. Finally, we note that the underlying discrete symmetry is free from the gravity argument against the axion.

Acknowledgments

I thank Bumseok Kyae for useful discussions. This work is supported in part by the National Research Foundation (NRF) grant funded by the Korean Government (MEST) (No. 2005-0093841).

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- [1] S. Weinberg, Phys. Rev. Lett. **19**, 1264 (1967).
[2] The ATLAS Collaboration, ‘*Observation of a new particle in the search for the Standard Model Higgs boson with the ATLAS detector at the LHC*,’ Phys. Lett. **B716**, 1 (2012) [arXiv:1207.7214];
The CMS Collaboration, ‘*Observation of a new boson at a mass of 125 GeV with the CMS experiment at the LHC*,’ Phys. Lett. **B716**, 30 (2012) [arXiv:1207.7235].
[3] J. E. Kim and H. P. Nilles, ‘*The μ -problem and the strong CP problem*,’ Phys. Lett. **B138**, 150 (1984).
[4] R. D. Peccei and H. R. Quinn, ‘*CP conservation in the presence of instantons*,’ Phys. Rev. Lett. **38**, 1440 (1977).
[5] G.F. Giudice and A. Masiero, ‘*A natural solution to the μ problem in supergravity theories*,’ Phys. Lett. **B206**, 480 (1988).
[6] J. A. Casas and C. Munoz, ‘*A natural solution to the μ problem*,’ Phys. Lett. **B306**, 288 (1993) [hep-ph/9302227].
[7] H. M. Lee, S. Raby, M. Ratz, G. R. Ross, R. Schieren, K. Schmidt-Hoberg, and P. K. S. Vaudrevange, ‘*A unique Z_4^R symmetry for the MSSM*,’ Phys. Lett. **B694**, 491 (2011) [arXiv:1009.0905].
[8] For a recent review, see, J. E. Kim and G. Carosi, ‘*Axions and the strong CP problem*,’ Rev. Mod. Phys. **82**, 557 (2010) [arXiv: 0807.3125[hep-ph]].
[9] J. E. Kim, ‘*Weak interaction singlet and strong CP invariance*,’ Phys. Rev. Lett. **43**, 103 (1979);
M. Dine, W. Fischler and M. Srednicki, ‘*A simple solution to the strong CP problem with a harmless axion*,’ Phys. Lett. **B104**, 199 (1961);
M. A. Shifman, A. I. Vainshtein and V. I. Zakharov, ‘*Can confinement ensure natural CP invariance of strong interactions?*,’ Nucl. Phys. **B166**, 493 (1980);
A. P. Zhitnitsky, ‘*On Possible Suppression of the Axion*

- Hadron Interactions*, Sov. J. Nucl. Phys. **31**, 260 (1980) [Yad. Fiz. **31**, 497 (1980)].
- [10] S. M. Barr and D. Seckel, ‘*Planck scale corrections to axion models*,’ Phys. Rev. **D46**, 539 (1992); M. Kamionkowski and J. March-Russell, ‘*Planck scale physics and the Peccei-Quinn mechanism*,’ Phys. Lett. **B282**, 137 (1992) [hep-th/9202003]; R. Holman, S. D. H. Hsu, T. W. Kephart, E. W. Kolb, R. Watkins, and L. M. Widrow, ‘*Solutions to the strong CP problem in a world with gravity*,’ Phys. Lett. **B282**, 132 (1992) [hep-ph/9203206]; S. Ghigna, M. Lusignoli and M. Roncadelli, ‘*Instability of the invisible axion*,’ Phys. Lett. **B283**, 278 (1992); B. A. Dobrescu, ‘*The strong CP problem versus Planck scale physics*,’ Phys. Rev. **D55**, 5826 (1997) [hep-ph/9609221].
- [11] J. E. Kim, ‘*Inverted-effective SUSY with combined Z' and gravity mediation, and muon anomalous magnetic moment*,’ Phys. Rev. **D87**, 015004 (2013) [arXiv:1208.5484].
- [12] S. Dimopoulos, S. Raby and F. Wilczek, ‘*Supersymmetry and the scale of unification*,’ Phys. Rev. **D24**, 1681 (1981); U. Amaldi, W. de Boer and H. Furstenuau, ‘*Comparison of grand unified theories with electroweak and strong coupling constants measured at LEP*,’ Phys. Lett. **B260**, 447 (1991); J. R. Ellis, G. Ridolfi and F. Zwirner, ‘*Radiative corrections to the masses of supersymmetric Higgs bosons*,’ Phys. Lett. **B257**, 83 (1991); P. Langacker and M.X. Luo, ‘*Implications of precision electroweak experiments for $M_t, \rho_0, \sin^2 \theta_W$ and grand unification*,’ Phys. Rev. **D44**, 817 (1991); C. Giunti, C. W. Kim and U. W. Lee, ‘*Running coupling constants and grand unification models*,’ Mod. Phys. Lett. **A6**, 1745 (1991).
- [13] J. E. Kim and M.-S. Seo, ‘*Mixing of axino and goldstino, and axino mass*,’ Nucl. Phys. **B864**, 296 (2012) [arXiv:1204.5495 [hep-ph]].
- [14] H. Harari, H. Haut and J. Weyers, ‘*Quark masses and Cabibbo angles*,’ Phys. Lett. **B78**, 459 (1978); Y. Chikashige, G. Gelmini, R. D. Peccei, and M. Roncadelli, ‘*Horizontal symmetries, dynamical symmetry breaking and neutrino masses*,’ Phys. Lett. **B94**, 499 (1980); P. Kaus and S. Meshkov, ‘*A BCS quark mass matrix*,’ Mod. Phys. Lett. **A3**, 1251 (1988), *ibid.* **A4**, 603 (1989)(E); H. Fritzsch and J. Plankl, ‘*Flavor democracy and the lepton-quark hierarchy*,’ Phys. Lett. **B237**, 451 (1990).
- [15] V. Braun, Y.-H. He, B. A. Ovrut, and T. Pantev, ‘*The exact MSSM spectrum from string theory*,’ JHEP **0605**, 043 (2006) [hep-th/0512177].
- [16] J. E. Kim, J.-H. Kim and B. Kyae, ‘*Superstring standard model from Z_{12-I} orbifold compactification with and without exotics, and effective R -parity*,’ JHEP **0706**, 034 (2007) [hep-ph/0702278].
- [17] J. E. Kim, ‘*A common scale for the invisible axion, local SUSY GUTs and saxino decay*,’ Phys. Lett. **B136**, 378 (1984).
- [18] For a recent SU(6), see, K.-S. Choi and J. E. Kim, ‘*Supersymmetric three family chiral SU(6) grand unification model from F-theory*,’ Phys. Rev. **D83**, 065016 (2011) [arXiv: 1012.0847[hep-ph]].
- [19] J. E. Kim, ‘*Axion and almost massless quark as ingredients of quintessence*,’ JHEP **9905**, 022 (1999) [hep-th/9811509]; ‘*Model-dependent axion as quintessence with almost massless hidden sector quarks*,’ *ibid.* **0006**, 016 (2000) [hep-ph/9907528]; K. Choi, ‘*Axions and the strong CP problem in M theory*,’ Phys. Rev. **D56**, 6588 (1997) [hep-th/9706171].
- [20] I.-W. Kim and J. E. Kim, ‘*Modification of decay constants of superstring axions: Effects of flux compactification and axion mixing*,’ Phys. Lett. **B639**, 342 (2006) [hep-th/0605256]; K.-S. Choi, H. P. Nilles, S. Ramos-Sanchez, and P. K. S. Vaudrevange, ‘*Accions*,’ Phys. Lett. **B675**, 381 (2009) [arXiv:0902.3070 [hep-th]]; J. E. Kim and H. P. Nilles, ‘*Axionic dark energy and a composite QCD axion*,’ JCAP **0905**, 010 (2009) [arXiv: 0902.3610 [hep-th]].
- [21] S. M. Barr, ‘*A new symmetry breaking pattern for SO(10) and proton decay*,’ Phys. Lett. **B112**, 219 (1982); J.-P. Derendinger, J. E. Kim and D. V. Nanopoulos, ‘*Anti-SU(5)*,’ Phys. Lett. **B139**, 170 (1984).
- [22] J. E. Kim and B. Kyae, ‘*Flipped SU(5) from Z_{12-I} orbifold with Wilson line*,’ Nucl. Phys. **B770**, 47 (2007) [arXiv:hep-th/0608086].
- [23] J.-H. Huh, J. E. Kim and B. Kyae, ‘ *$SU(5)_{flip} \times SU(5)'$ from Z_{12-I}* ,’ Phys. Rev. **D80**, 115012 (2009) [arXiv:0904.1108].
- [24] I. Antoniadis, J. R. Ellis, J. S. Hagelin, and D. V. Nanopoulos, ‘*GUT model building with fermionic four-dimensional strings*,’ Phys. Lett. **B205**, 459 (1988).